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Argonne National Laboratory

PHYSICS DIVISION SUMMARY REPORT

September-December 1963

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ARGONNE NATIONAL LABORATORY
9700 South Cass Avenue
Argonne, Illinois 60440

PHYSICS DIVISION
SUMMARY REPORT

September-December 1963

Lowell M. Bollinger, Division Director

Preceding Summary Reports:

ANL-6719, March-May 1963
ANL-6720, June 1963
ANL-6767, July-August 1963

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FOREWORD

This issue of the Argonne National Laboratory Physics Division Summary Report has been considerably delayed—first to prepare the proceedings of the International Conference on Nuclear Physics with Reactor Neutrons and soon thereafter to prepare the contributed papers and the proceedings of the Symposium on Nuclear Spectroscopy with Direct Reactions.

The articles in the Physics Summary are informal progress reports. The results and data therefore must be understood to be preliminary, tentative, and often incomplete.

The issuance of these reports is not intended to constitute publication in any sense of the word. Final results either will be submitted for publication in regular professional journals or, in special cases, will be presented in ANL Topical Reports.

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I. EXPERIMENTAL NUCLEAR PHYSICS

I-80-33. MOLECULAR-BEAM STUDIES* (51210-01)

William Childs, John Dalman, and Leonard Goodman

Since the last summary report, a great deal of new data has been taken with the Mark II atomic-beam machine. These results may be loosely separated into two groups: (1) precision measurement of electronic g factors of metastable atomic states, and (2) measurement of the magnetic-dipole hyperfine-interaction constant $\underline{a}$ for various atomic states in different isotopes. The information on g_J values is summarized in Tables I and II, and that on hyperfine-interaction constants in Table III.

TABLE I. Measured values of electronic g factors of metastable atomic states in Fe, Ge, and Sn. The figures in parentheses indicate the uncertainty in the last decimal place.

Element	Metastable atomic state	Excitation energy of state (cm^{-1})	Measured g factor
Fe	$^5F_5$	6928	1.40007(6)
	$^5F_4$	7377	1.34995(7)
	$^5F_3$	7728	1.254(7)
Ge	$^3P_1$	557	1.50111(7)
	$^3P_2$	1410	1.49458(9)
	$^1D_2$	7125	1.00639(8)
Sn	$^3P_1$	1692	1.50110(7)
	$^3P_2$	3428	1.44878(9)
	$^1D_2$	8613	1.05230(8)

*This report was written in November 1963 and represents research

TABLE II. Departure of the g factors g_{exp} of the $p^2(^1D_2, ^3P_{1,2})$ states from the values g_{RS} predicted on the basis of Russell-Saunders coupling. The figures in parentheses indicate the uncertainty in the last decimal place. The equal and opposite perturbation of the g factors of the J=2 states is clear.

Atomic state	$g_{exp} - g_{RS}$	
	Ge I	Sn I
$^3P_1$	-0.00005(7)	-0.00006(7)
$^3P_2$	-0.00658(9)	-0.05238(9)
$^1D_2$	+0.00639(8)	+0.05230(8)

TABLE III. Measured values of the magnetic hyperfine-interaction constant a for three nuclei. The value of the magnetic field H_J produced at the nucleus by the electronic configuration is believed to be accurate to within about 1%.

Isotope	Atomic state	Excitation energy of state (cm^{-1})	$\|a\|$ (Mc/sec)	$\|H_J\|$ (Gauss)
Fe^{57}	$^5D_4$	Ground state	38.14 ± 0.28	1.10×10^6
Sn^{117}	$^3P_1$	1692	553.0 ± 0.5	3.63×10^5
	$^3P_2$	3428	1221 ± 6	1.60×10^6
Sn^{119}	$^3P_1$	1692	578.8 ± 0.5	3.63×10^5
	$^3P_2$	3428	1272 ± 9	1.60×10^6

Prior to the present work, no atomic-beam investigations of metastable states lying as high as 6000 cm^{-1} were possible without the use of special techniques for excitation of the beam atoms. For this reason, we were somewhat surprised to find resonances from such states interfering with our work on lower states. It is the excellent signal-to-noise ratio of the Mark II machine that makes such investigations possible.

The g factors given for Fe were obtained relative to the precise value previously obtained for the ^5D ground state,[1] and are not to be considered as final values.

A paper describing the work on the g factors of Ge and Sn has been written for the Physical Review. Both elements have the electronic configuration ns^2np^2, which leads to the terms $^1S_0$, $^1D_2$, and $^3P_{0,1,2}$, of which the $^1D_2$ and $^3P_{1,2}$ states have measurable g factors. The energy spacings of these states show evidence of a considerable amount of spin-orbit mixing, particularly in Sn. Because the spin-orbit interaction can mix only states of the same J, the g value of the $^3P_1$ state should and does lie extremely close to the Russell-Saunders value of 1.50116 for both Ge and Sn. The $^1D_2$ and $^3P_2$ states can, however, perturb each other strongly. Table II shows the difference between the measured values and the Russell-Saunders values of the g factors for the three states. It is seen that, within experimental error, the g factors of the $^1D_2$ and $^3P_2$ states are shifted by equal and opposite amounts from the Russell-Saunders limits in both Ge and Sn. An intermediate-coupling calculation accounts for 96% of the shift in Ge and 85% of that in Sn. The remaining discrepancy is attributed to configuration interaction.

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[1] A. Kamlah and S. Penselin (private communication, 1962).

From the value of the hyperfine-interaction constant a given for the $^5D_4$ atomic ground state of Fe^{57}, the value $(1.10 \pm 1\%) \times 10^6$ G can be calculated for the magnetic field at the iron nucleus in the atomic ground state of the free atom.

The measurements on Sn^{117} and Sn^{119} in the metastable $^3P_1$ and $^3P_2$ states are summarized in Table III. Observation of the $\Delta F = \pm 1$ transitions must be postponed until suitable rf signal generators can be obtained. The results given for the hyperfine-interaction constants in Fe^{57} and $Sn^{117,119}$ are not to be considered as final values.

Note: In the considerable period since this report was written, two papers presenting final results have been published. These are (1) "Electronic g Factors of the p^2 Configurations in Ge I and Sn I" [W. J. Childs and L. S. Goodman, Phys. Rev. 134, A66 (1964)] and (2) "Magnetic Hyperfine Structure of the $^3P_1$ and $^3P_2$ Metastable States of $Sn^{115,117,119}$" [W. J. Childs and L. S. Goodman, Phys. Rev. 137, A35 (1965)].

II. MASS SPECTROSCOPY

II-40-12. FRAGMENTATION OF HYDROCARBONS (51300-01)

H. E. Stanton and J. E. Monahan

The method of data analysis proposed in an earlier report[1] has been subjected to considerable development and testing. The method comprises computing the Laplace transform of an unknown kinetic energy distribution as measured by mass spectrometer MA-17, repeating this for a known standard, and finding the ratio of the two transforms. From this ratio one can deduce some parameters of the unknown distribution. In principle, one can even invert the transformation and exhibit the kinetic energy distribution of the ion freed from the effects of measurements, i.e., freed from the smoothing due to imperfect energy resolution, the spread in the electron beam, etc. This complete inversion of the Laplace transforms poses certain difficulties, however. In particular, statistical fluctuations in the form of noise mean that the data are not "perfect," and it seemed worthwhile to try an approach which was less exact and exacting.

A Laplace transform is defined by the expression

$$\tilde{Y}(S) = \int_0^\infty e^{-SE}\, Y(E)\, dE, \qquad (1)$$

where $\tilde{Y}(S)$ and $Y(E)$ are transform pairs. From this it is clear that

$$\tilde{Y}(0) = \int_0^\infty Y(E)\, dE, \qquad (2)$$

[1] H. E. Stanton and J. E. Monahan, Physics Division Summary Report ANL-6679, January-February 1963, p. 56.

$$\lim_{S\to 0} \frac{d}{dS} \tilde{Y}(S) = - \int_0^\infty E\, Y(E)\, dE, \tag{3}$$

$$\lim_{S\to 0} \frac{d^n}{dS^n} \tilde{Y}(S) = (-1)^n \int_0^\infty E^n\, Y(E)\, dE, \tag{4}$$

provided certain mathematical restrictions (such as proper behavior of the integrals) are satisfied. If Y(E) is interpreted as an energy distribution curve, Eq. (2) represents the total number of particles, Eq. (3) represents the first moment of the distribution, and Eq. (4) represents the general nth moment of the distribution.

Let $M_k(E)$ and $M_u(E)$ be the actual measured distributions in kinetic energy for the standard ion and the unknown ion, respectively, while $Y_k(E)$ and $Y_u(E)$ represent the actual distributions when the effects of instrumentation are eliminated. As shown in Ref. 1, the ratio

$$R(S) = \frac{\tilde{Y}_k(0)}{\tilde{Y}_k(S)} \cdot \frac{\tilde{Y}_u(S)}{\tilde{Y}_u(0)} \tag{5}$$

is independent of the effects of instrumentation.

If the total ion-current distribution, recorded separately on MA-17, is used to match the energy scales for M_k and M_u, there may be errors in energy displacements $\mathcal{E}_{0u}$ for the unknown and $\mathcal{E}_{0k}$ for the standard. Under these conditions[1]

$$\frac{\tilde{M}_u(S)}{\tilde{M}_u(0)} = \frac{\tilde{Y}_u(S)}{\tilde{Y}_u(0)} \exp(S\mathcal{E}_{0u}) . \tag{6}$$

Hence

$$\frac{1}{\tilde{M}_u(0)} \left(\frac{d\tilde{M}_u(S)}{dS}\right)_0 = \mathcal{E}_{0u} + \frac{1}{\tilde{Y}_u(0)} \left(\frac{d\tilde{Y}_u(S)}{dS}\right)_0 , \tag{7}$$

and similarly for the standard.

One may now differentiate Eq. (5). Using Eq. (7) and the analogous expression for the standard ion and taking the limit as $S \to 0$ leads to

$$\frac{\tilde{Y}_u'(0)}{\tilde{Y}_u(0)} - \frac{\tilde{Y}_k'(0)}{\tilde{Y}_k(0)} = \frac{\tilde{M}_u'(0)}{\tilde{M}_u(0)} - \frac{\tilde{M}_k'(0)}{\tilde{M}_k(0)} - \mathcal{E}_{0u} + \mathcal{E}_{0p}. \quad (8)$$

With the aid of Eqs. (2) and (3), this expression relates the moments of the intrinsic distribution ($\tilde{Y}$) of the ions to the moments readily obtainable from the data ($\tilde{M}$) provided the values of $\mathcal{E}_{0u}$ and $\mathcal{E}_{0p}$ are known. These were evaluated by use of the trace for the total ion beam measured after energy analysis in the mass spectrometer. Two additional relations similar to Eq. (8) were derived for the two next higher moments of the distributions by successive differentiations of Eq. (5). The derivations were lengthy and involved and the results will not be shown.

In order to relate these developments to the intrinsic kinetic energy distributions in the fragmentation process, it is necessary to transform the fragmentation-energy distribution in the center-of-mass system of the parent molecule to the laboratory system.[2] Let the distribution of the fragmentation energy Q of a particular ionic fragment be denoted by A(Q). The transformation can be carried out by standard means. If the parent molecules are moving according to a Maxwell-Boltzmann distribution of energies, then

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[2] H. E. Stanton and J. E. Monahan, Physics Division Summary Report ANL-6326, March 1961, p. 55.

$$Y(E) = \left(\frac{1}{2\pi\theta}\,\frac{M}{m_1 m_2}\right)^{1/2} \int_0^\infty dQ\,\frac{A(Q)}{Q}\left\{\exp\frac{-[(ME)^{1/2} - (m_2Q)^{1/2}]^2}{m_1\theta}\right.$$

$$\left. - \exp\frac{-[(ME)^{1/2} + (m_2Q)^{1/2}]^2}{m_1\theta}\right\}, \qquad (9)$$

where θ is the temperature and M, m_1, and m_2 are the masses of the parent molecule ion, the ionic fragment, and the neutral or charged companion. The Laplace transform of Eq. (9) is

$$\tilde{Y}(S) = \frac{1}{(m_1/M)\theta S + 1} \int_0^\infty dQ\; A(Q)\; \exp\frac{-(m_2/M)SQ}{(m_1/M)\theta S + 1}\,. \qquad (10)$$

The properties of the distribution A(Q) that can be obtained from this analysis consist in the moments

$$P_1 = \int_0^\infty dQ\; Q\; A(Q), \qquad (11)$$

$$P_i = \int_0^\infty dQ\; (Q - P_1)^i\; A(Q)\,.$$

These are evaluated by computations from Eqs. (4), (5), and (8) and related equations for higher moments. This implies that two sets of moments are involved in the calculations since R(S) is the ratio of the normalized transform for the standard to that for the unknown. The derivations are greatly simplified if one chooses a parent-ion distribution as the standard because then $A(Q) = \delta(Q)$ and

$$P_{1k} = \mathcal{E}_{0k}\,,$$

$$P_{ik} = 0 \quad \text{for} \quad i > 1\,.$$

Consequently, except for the first moment, it follows that if P_i represents the difference between the moments of the unknown distribution and the parent ion distribution, then

$$P_1 = \frac{M}{m_2} R'(0) - \frac{3}{2}\theta,$$

$$P_2 = \frac{M^2}{m_1^2}\{R''(0) - [R'(0)]^2\} + \frac{m_2}{m_1}\theta\left[\frac{3}{2}\theta\left(\frac{M}{m_1} + 3\right) - 2\frac{M}{m_2}R'(0)\right], \quad (12)$$

where

$$R''(0) = \lim_{S\to 0} \frac{d^2}{dS^2} R(S),$$

etc. Analogous expressions for P_3, etc. are too complicated to include. In any event, a sequence of internal moments of the fragmentation distribution A(Q) is directly derivable from the experimental data.

In the following analysis, only the first three moments were computed and used. Two possible types of distributions have been tested as follows. The first assumption was that

$$A(Q) = A_n Q^n e^{-Q/Q_0}, \quad (13)$$

where

$$A_n = \frac{1}{Q_0^{n+1}\Gamma(n+1)}$$

and therefore

$$\int_0^\infty dQ\, A(Q) = 1.$$

The latter represents a generalized Maxwell distribution. From this, one can compute the moments

$$P_1 = (n + 1)\, Q_0 \,,$$
$$P_2 = (n + 1)\, Q_0^{\,2} \,, \tag{14}$$
$$P_3 = 2(n + 1)\, Q_0^{\,3} \,.$$

It is clear that these three moments must satisfy the expression

$$P_1 P_3 = 2P_2^{\,2} \tag{15}$$

in order for the distribution to be Maxwellian.

The second assumption was that the distribution A(Q) consisted of two sharp distributions with energies q_1 and q_2 and amplitudes $(\frac{1}{2} - a)$ and $(\frac{1}{2} + a)$, respectively; i.e.,

$$A(Q) = (\tfrac{1}{2} - a)\, \delta(Q - q_1) + (\tfrac{1}{2} + a)\, \delta(Q - q_2) \,, \tag{16}$$

where

$$\int_0^\infty dQ\, A(Q) = 1,$$

and the parameters a, q_1, and q_2 are to be determined. This yields the moments

$$P_1 = (\tfrac{1}{2} - a) q_1 + (\tfrac{1}{2} + a) q_2 \,,$$
$$P_2 = (\tfrac{1}{4} - a^2)(q_2 - q_1)^2 \,, \tag{17}$$
$$P_3 = 2a(\tfrac{1}{4} - a^2)(q_1 - q_2)^3 \,,$$

which may be used to determine the parameters a, q_1, and q_2 of the assumed distributions.

As a special case of this for a parent peak for which there is no fragmentation energy, it turns out that $a = \pm \frac{1}{2}$ only. When studying

ionic distributions in this class, one can only measure the temperature of the target gas; and this temperature can be used to check an earlier theory of distribution analysis[2] and instrumental performance.

Experimentally, these theoretical developments were tested by the analysis of three molecules: H_2O, CH_4, and CO. The latter was used to compare with results previously found by Hagstrum[3] for O^+ and C^+. Before taking up the analysis of these molecules, it seems advisable to consider any possible mass discriminations in the instrument.

The ability to discriminate between ions of different energies was determined by the widths S_1 and S_2 of the entrance and exit slits of the energy analyzer, the radius of the analyzer (R = 12 in.), and the initial accelerating voltage E_0. These constants define the quantity $b = (S_1 + S_2)E_0/2R$ which is the approximate range of energies transmitted for a given setting of the controls. The value of b was usually 0.17 eV or less. This corresponded to a value E_0 = 400 V which represented a compromise between the desire for a good signal-to-noise ratio for small peaks and the conflicting desire for simplicity in the energy analysis. In practice, the effects of focusing would at times further reduce the value of b appreciably.

EXPERIMENTAL RESULTS

The measured distribution in ion kinetic energy depends upon conditions in the ion chamber in an important way. The total kinetic energy acquired by an ion at any given point along its trajectory will be made up of any kinetic energy of formation and, in addition, the difference in electrostatic potential between the given point and the

[3] H. D. Hagstrum, Rev. Mod. Phys. 23, 185 (1951).

point of formation. Consequently, even though formed at rest, ions will show a kinetic energy spread because they are formed at different depths in the ion chamber where they are accelerated over different distances in the potential gradients imposed by the necessary drawing-out potentials. The errors introduced in this way may be an order of magnitude larger than any kinetic energy of formation. In the present investigation ions were formed between two plates, parallel to the electron trajectory and at right angles to the exit direction for the ions. Controllable potential gradients along the ion axis in the ion chamber could be produced. It was found that an effective spread of $\lesssim 0.007$ in. in the electron beam produced a kinetic energy distribution of ions from argon gas with a width at half maximum of ≈ 0.5 eV including the energy spread 2b. Usually, however, this width is less than 0.25 eV and has been as low as 0.15 eV. Frequently surface conditions in the ion chamber may distort and defocus the electron beam.

The techniques of these experiments and their interpretation involve the comparison of an unknown distribution with a known (standard) one for which the P_i can be computed. Parent ions, including the rare gas ions, are suitable for this purpose since they do not acquire translational energy as a result of electron impact and may be presumed to be at thermal equilibrium with the ion chamber.

The variation with mass number may be measured by the intercomparison of "parent" ion distributions [e.g., A^{+++} ($m/e = 13\frac{1}{3}$), H_2O^+, etc.] which were used experimentally. Since, however, the comparison involves the arithmetic difference between corresponding P's of two distributions, these differences should theoretically be zero if the peaks are measured all at the same temperature and provided there is no mass dependence of either the P's themselves or their differences.

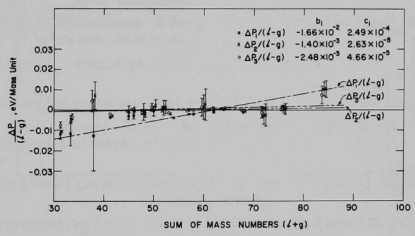

Fig. 1. The behavior of the first three moments of the center-of-mass distributions for the parent ions CO_3^+, A^+, O_2^+, CO^+, A^{++}, H_2O^+, and A^{+++}. Ordinates are differences in moments divided by differences in mass for pairs of ions; abscissas are sums of masses for the same pairs of ions. The slope and intercept give the parameters in the expression for mass dependence.

If the subscripts i and j in the expression $Q_{ij} = P_k$ for the kth moment refer to the ionic mass numbers of the ith ion as standard and the jth parent ion, then the function empirically relating the moments to the masses may be assumed to be expressible in the form

$$Q_{ij} = a_i + jb_i + j^2 c_i \,. \tag{18}$$

One notes that a_i cannot be determined if only differences in Q's between pairs of peaks are substituted in this equation. If, next, each ion is used as the standard in sequence, six ions can supply a large number of points.

The kinetic energy distribution of each of the ions A^{+3}, H_2O^+, A^{+2}, N_2^+ (CO^+), O_2^+, A^+, and CO_2^+ was used in turn as a standard, and the measurements of the first three moments computed from the other distributions were compared with it. The resulting points are plotted in Fig. 1 together with the probable errors of the P's for each value of mass, and the straight lines which best fit the points.

These data have been analyzed in the following way. It follows from Eq. (18) that for two masses

$$\begin{aligned} \Delta Q_{i\ell g} &= Q_{i\ell} - Q_{ig} \\ &= (\ell - g)b_i + (\ell^2 - g^2)c_i \,, \\ \frac{\Delta Q_{i\ell g}}{\ell - g} &= b_i + (\ell + g)c_i \,. \end{aligned} \tag{19}$$

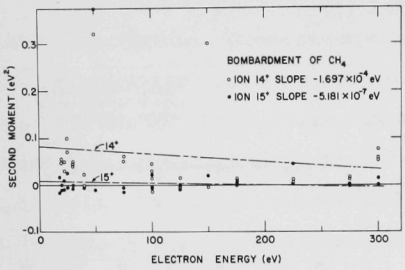

Fig. 2. Sample plot of moment vs electron energy for a particular fragmentation process. In this instance, the variations of the second moment for electrons with energies from 18 eV to 300 eV are shown for the CH_3^+ and CH_2^+ fragments of methane. Theoretical reasons for a negative slope are not immediately apparent.

Consequently a plot of $\Delta Q/(\ell - g)$ for the various combinations is linear in the sum of the mass numbers to this degree of approximation. This is shown in Fig. 1 for all three moments together with a least-squares fit for each. According to Eq. (11), P_1 depends on energy in a way which is different from that for the other moments. In order to calculate P_1, the distribution for the total ion beam issuing from the energy analyzer was used. Unfortunately the different locations of the two detectors permitted stray magnetic fields to affect their readings in different ways. Lack of mechanical rigidity could also produce spurious displacements of the two distributions and thus cause a higher variation in P_1 than for P_2 and P_3. The latter are moments intrinsic to the distribution of a particular ion and the lack of much variation with mass number would seem to imply that the shapes of the distributions were reasonably well preserved with changes in mass.

In analyzing the experimental distributions in terms of the Boltzmann distribution of Eq. (13) with three moments computed for the center-of-mass distribution, one can find the exponent n and "temperature" Q_0 from any pair of moments and use the third to examine consistency.

Figure 2 illustrates the results of calculating the second moments of the fragmentation distribution in energy from the measured kinetic energy distributions. The variable in this case was the electron

bombarding energy. The quantity displayed is the second moment for the ions CH_3^+ and CH_2^+ from methane. Similar plots were made for all fragments in the molecules examined, except CO^+ and those containing only hydrogen, and for the other moments derived therefrom. In some cases the scatter of the points was unfortunately large. In Fig. 2, for instance, all but one of the values at an electron energy of 50 eV appear to be too far off to be meaningful, but were included in determining the line that best fitted the data.

Values of the moments used for the calculations were obtained from a least-squares fit of a series of observations at different energies of the bombarding electrons. Three values were selected, corresponding to electron energies of "0 eV" (extrapolation to zero-energy electrons), 150 eV, and 300 eV. The parameters were then determined from one pair of moments for each of these three bombarding energies and averaged. The results are listed in Table IV. Three values appear at each position in the table since there were three possible pairs of moments to determine the two parameters.

It is apparent from the table that the value of the exponent n is by no means uniquely determined except for the two parent peaks (CH_4^+ and H_2O^+) and possibly OH^+. The values of Q_0 also showed wide variations and where blanks appear in the table, were negative. The lack of uniqueness of n and the erratic behavior of Q_0 for the parent peaks can be explained by the fact that the parent ion can have no energy of fragmentation in the center-of-mass system; the behavior of the parameters results from trying to fit a delta function at $Q = 0$ with a Boltzmann distribution.

Also, with the other fragments, a Boltzmann distribution cannot be made to accommodate the real distribution in the center-of-mass system. One notes, however, that two values of n are usually close and negative, while the third is large and positive. The former involve the

TABLE IV. The ions studied are listed in column 1, the values of the exponent n in the Boltzmann distribution are listed next with three entries for each ion depending on which pair of the P's were used. The first entry is the value of n derived from P_1 and P_2, the second from P_1 and P_3, and the third from P_2 and P_3. The values of Q_0 [Eq. (14)] are listed in the third column, again with three entries computed with the same three combinations of P's.

Fragment	Values of n	Values of Q_0
	Methane	
CH_4^+	-0.982, -0.989, -0.973	· · · , 0.406, · · ·
CH_3^+	-0.759, -0.851, 3.46	0.116, 0.205, 0.313
CH_2^+	-0.892, -0.731, 4.92	1.014, 0.302, 0.096
CH^+	-0.602, -0.648, 1.00	0.583, 0.540, 1.44
C^+	-0.775, -0.455, 6.82	0.945, 0.390, 0.155
	Water	
H_2O^+	-1.00 , -1.00 , -1.00	· · · , 1.99 , 0.470
OH^+	-0.943, -0.945, -0.887	0.857, 0.636, 0.520
O^+	-0.406, 0.140, 7.56	0.414, 0.212, 0.109
	Carbon Monoxide	
O^+	0.140, 0.720, 4.90	1.50 , 0.984, 0.666
C^+	-0.512, -0.278, 1.37	· · · , 0.463, 0.393

first moment in the calculations whereas the latter is calculated from the other two only.

The second method, described in Eq. (16) et seq., was also used. Here again, there were troubles which are attributed to the data, particularly the first moment. Values of q_1 in Eq. (16) would frequently turn out to be negative. This was attributed to P_1 in Eq. (17). Usually a slight readjustment of its value would make $q_1 \geq 0$. The imaginary values occasionally obtained for q_1 or q_2 indicate either that the data are of poor quality or that they cannot be fitted with the distribution assumed.

The results of this analysis are shown by the bars in Fig. 2, where average values for q_1 and q_2 were used. The height of a bar indicates its relative importance. In the case of the parent peaks, with practically no readjustment, there is a single sharp peak at zero energy and a negligibly small second peak erratically placed at a relatively high value of energy (either positive or negative). Since for parent peaks there should be a single sharp energy of fragmentation (at zero), it seems reasonable to assume that the general method of analysis has some validity. Values of q_1 and q_2 and their amplitudes are given in Table V.

Figure 3 shows the Boltzmann distributions calculated with appropriate values of n and Q_0 taken from Table IV and with arbitrary amplitudes. In all cases except H_2O^+, OH^+, and O^+ from CO, one negative value and one positive value of the exponent n was used. Usually an average Q_0 was used. For O^+ from CO, two positive values were used for n.

As one might expect, the values of q_1 and q_2 seem to determine the approximate values of the first moments of these calculated distributions. This would seem to follow from the fact that the three moments plus a normalization condition on the Boltzmann

TABLE V. Amplitudes and energies of the two sharp components assumed in the center-of-mass distributions in fragment kinetic energy. The ions investigated are listed in column 1 for each of the three vapors used.

Fragment	Low-energy component		High-energy component	
	Energy (eV)	Amplitude	Energy (eV)	Amplitude
		Methane		
CH_4^+	0	1.00	. . .	. . .
CH_3^+	0.035	0.984	0.654	0.016
CH_2^+	0.007	0.947	1.18	0.053
CH^+	0.021	0.824	1.20	0.176
C^+	0.020	0.843	1.33	0.157
		Water		
H_2O^+	0.0015	1.00	6.77	3.9×10^{-5}
OH^+	0.015	0.985	1.29	0.015
O^+	0.011	0.664	0.693	0.336
		Carbon Monoxide		
C^+	0.061	0.773	0.982	0.227
O^+	0.638	0.697	4.09	0.301

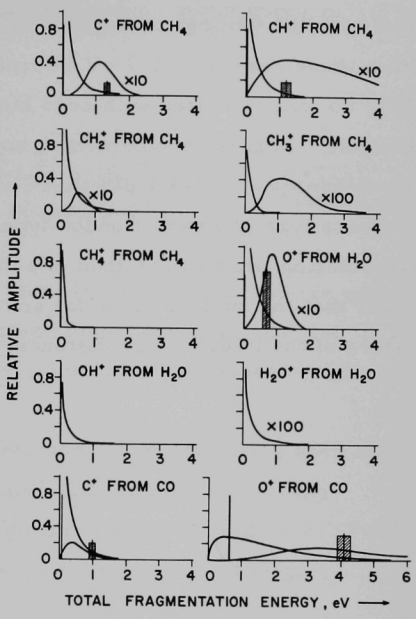

Fig. 3. The two types of distribution used in this analysis, plotted as functions of the total kinetic energy of fragmentation. The energies q_1 and q_2 of the two sharp center-of-mass distributions assumed were determined by fitting the measured distributions. The low-energy component is shown as a narrow line close to the origin, the high-energy one as a shaded bar (whose width indicates the error). The fragments CH_4^+, H_2O^+, and OH^+ showed no high-energy components. The curves represent the two Boltzmann distributions ascribable to all peaks, except those mentioned above. The two separate curves for each ion are for the negative and the positive value of the exponent n. The area enclosed by each of these curves was adjusted to be proportional to the amplitude of the corresponding sharp distribution.

distributions permit fitting the experimental distribution with two Boltzmann distributions. If now these distributions are normalized to the amplitudes associated with q_1 and q_2, one might expect their first moments to agree with these two fragment energies.

Unfortunately these analyses do not seem to justify much inference concerning the fundamental fragmentation process. It seems rather clear that a "boiling off" process of fragmentation (which one would expect to give some sort of Boltzmann distribution) does not, in fact, agree with these results. In the case of methane fragmentation, the relative amplitudes of q_1 and q_2 imply that a large portion of the ions are liberated with zero energy in the center-of-mass system as one would expect from the "equilibrium theory" of fragmentation.

The distribution in energy of O^+ from CO agrees reasonably well with the findings of Hagstrum. The average kinetic energy derived from these results is 1.7 eV whereas he finds about 2.0 eV. He also finds no ions liberated with zero kinetic energy, in agreement with this analysis. In these experiments, however, we find many more ions of lower energy (i.e., <1 eV) than Hagstrum. This could indicate an ordinary type of Frank-Condon type transition that includes the dissociation limit rather than one that does not. This would in turn imply that the nuclear potential energy curve is shifted to smaller values of the internuclear distance.

V. THEORETICAL PHYSICS, GENERAL

V-19-1. NOTE ON A CRUDE NUCLEON-NUCLEON POTENTIAL

F. Coester and Edward Yen* (51210-01)

For the investigation of nuclear-structure problems with "realistic" nuclear forces, it is desirable to have a simple analytic nucleon-nucleon potential which roughly reproduces the most important two-body scattering data.

Potentials which have been constructed to give a good fit to all available data are fairly complicated and the required strong repulsion at high energies is usually approximated by a hard core. For such potentials, two-body matrix elements with products of single-particle wave functions do not exist.

A useful crude potential should be mathematically simple, representing only the essential features of the nuclear force. Of course, which features are essential depends somewhat on the purpose at hand. For example, the tensor force is essential for some purposes while a purely central force may be preferred for the sake of simplicity where it is adequate.

For the present, spin-dependent central Serber potentials are being considered. A necessary condition to prevent collapse of heavy nuclei is

$$\int_0^\infty dr\, r^2\, V(r) \geq 0 \,. \qquad (1)$$

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We assume

$$\int_0^\infty dr\, r^2\, V(r) = 0 \,. \tag{2}$$

An example of such a potential is

$$V(r) = V_0\, (4\, e^{-2\mu r} - e^{-\mu r})/\mu r \,, \tag{3}$$

which is strongly repulsive at short distances but does not contain a hard core.

The strength and the range of this potential are determined to fit the neutron-proton scattering lengths and effective ranges,[1] which are $a_t = 5.396 \pm 0.011$ F, $a_s = -23.678 \pm 0.028$ F; $r_t = 1.726 \pm 0.014$ F; $r_s = 2.51 \pm 0.15$ F. The potential parameters are summarized in Table VI. Figure 4 shows the potential shape. The singlet potential is for $V_0 = 2307$ MeV.

After the strengths and ranges are determined by low-energy scattering data, the high-energy p-p singlet phase shifts for s and d waves have been calculated and compared with those of Breit et al.[2] These data are given in Table VII. We conclude that the s-wave

TABLE VI. Potential parameters, computed scattering length, and effective range.

S	V_0 (MeV)	μ/m_π	a (F)	r_0 (F)
0	2126	2.81	23.68	2.62
0	2307	2.92	23.68	2.51
1	3991	3.26	5.394	1.726

[1] H. P. Noyes, Phys. Rev. 130, 2025 (1963).

[2] G. Breit, M. H. Hull, K. E. Lassila, and K. D. Pyatt, Phys. Rev. 120, 2227 (1960).

interaction can be described adequately by the potential proposed in Eq. (3). A good fit was not to be expected for the d-wave phase shifts.

Since the range of the potential (3) is much shorter than that of the well-established one-pion-exchange potential, the same computations have been repeated for the potential

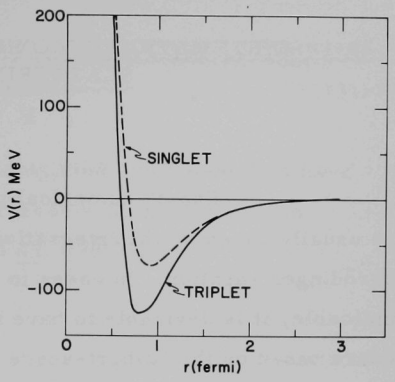

Fig. 4. Potential shape.

$$V(r) = V_0 \frac{(4\, e^{-2\mu r} - e^{-\mu r})}{\mu r} + V_\pi \frac{(4\, e^{-2m_\pi r} - e^{-m_\pi r})}{m_\pi r}, \qquad (4)$$

where $V_\pi = 11.17$ MeV and $m_\pi^{-1} = 1.413$ F. The high-energy s-wave phase shifts are too large in that case.

TABLE VII. Comparison of experimental and computed phase shifts.

E (MeV)	L	δ (YLAM)	δ	
			$V_0 = 2307$ MeV	$V_0 = 2126$ MeV
95	0	0.419	0.426	0.390
210	0	0.087	0.045	0.063
310	0	-0.175	-0.168	-0.212
95	2	0.068	0.081	0.09
210	2	0.140	0.22	0.23
310	2	0.209	0.29	0.30

V-20-1. SYSTEMATIC APPROXIMATIONS FOR THE SINGLE-CHANNEL SCATTERING AMPLITUDE[1] (51210-01)

F. Coester

"Exact" numerical computations of scattering amplitudes are usually based on the integration of a time-independent differential Schrödinger equation. In cases in which these methods are not applicable, it is desirable to have a systematic approximation procedure based on the Hilbert-space formulation of scattering theory.

The scattering amplitude for two-particle scattering may be computed by inverting a Hilbert-space operator (1 - K), where K is completely continuous and depends on the real energy as a parameter. This inversion can always be accomplished by obtaining a finite number of eigenvalues and eigenfunctions of K together with a modified Born expansion which is guaranteed to converge. Another method consists in the inversion of a finite matrix, the elements of which are computed by a convergent perturbation series. These methods are also applicable to coupled two-particle channels.

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[1] F. Coester, Phys. Rev. 133, B1516 (1964).

V-26-2. PREACCELERATION IN ELECTRON THEORY

M. N. Hack (51151-01)

This project has been completed and the results have been published in a report entitled "Preacceleration in Electron Theory," M. N. Hack, Nuovo Cimento 29, 298 (1963).

V-47-2. STRUCTURE OF ELEMENTARY PARTICLES

K. Hiida (51151-01)

PSEUDOSCALAR CHARGE DENSITY OF SPIN-$\frac{1}{2}$ PARTICLES WITH PARITY-NONCONSERVING INTERACTIONS

1. Renormalization

Consider the Lagrangian density for a spin-$\frac{1}{2}$ particle with mechanical mass m_0, namely

$$L = L_1 + L_2, \qquad L_1 = -\bar{\psi}(x)\left[\gamma_\mu \frac{\partial}{\partial x_\mu} + m_0\right]\psi(x), \tag{1}$$

where the Dirac gamma matrix γ_μ satisfies the commutation relation

$$\{\gamma_\mu, \gamma_\nu\} = 2\delta_{\mu\nu} \tag{2}$$

and L_2 is the interaction Lagrangian density, assumed to be CP invariant (but not P invariant). For definiteness, for a moment, we shall consider only renormalizable interactions. Since the term $\bar{\psi}\gamma_\mu\gamma_5(\partial/\partial x_\mu)\psi$ is invariant and the term $\bar{\psi}\gamma_5\psi$ is not invariant under CP transformation, the former term as well as the self-energy term $\bar{\psi}\psi$ should be induced by the interaction of the spin-$\frac{1}{2}$ particle with its self field, where $\gamma_5^2 = 1$.

Now our proposal is that the terms induced by the self interaction should be added to the free part of the Lagrangian density L_1 and the same terms be subtracted from the interaction Lagrangian density L_2, in order to remove all divergences from the renormalized S matrix. Then our Lagrangian density is given by

$$\begin{cases} L = L_0 + L', \\ L_0 = -\bar{\psi}(x)\left[\Gamma_\mu \dfrac{\partial}{\partial x_\mu} + m\right]\psi(x), \\ L' = L_2 + \delta m\, \bar{\psi}(x)\psi(x) + a(1-a^2)^{-1/2}\, \bar{\psi}(x)\gamma_\mu\gamma_5 \dfrac{\partial}{\partial x_\mu}\psi(x), \end{cases} \tag{3}$$

where a is a real constant under the requirement that L_2 is CP invariant, $a^2 < 1$ by the requirement that the particle cannot propagate in vacuum faster than the velocity of light, and the definition of Γ_μ is

$$\Gamma_\mu = \gamma_\mu(1 + a\gamma_5)(1 - a^2)^{-1/2}, \tag{4}$$

which satisfies the same commutation relation as γ_μ, namely

$$\{\Gamma_\mu, \Gamma_\nu\} = 2\delta_{\mu\nu}. \tag{5}$$

It should be noted that the definition of $\bar{\psi}$ in (3) is still $\psi^*\beta$ but not $\psi^*\beta(1 + a\gamma_5)(1 - a^2)^{-1/2}$, because the term $\gamma_\mu\gamma_5$ in Γ_μ comes from the self interaction.

Starting from the Lagrangian density (3), we find that the unrenormalized modified propagator of spin-$\frac{1}{2}$ particles with nonvanishing mass is expressed as

$$\begin{aligned} S_F^{\,1}(p) = i\Big\{ & -\delta m - \frac{a^2 m^2}{1 - a^2} + z\int_0^\infty dx^2\,[(m - x)\sigma_1(x^2) + \rho_2(x^2)]\Big\} S_F^{\,2}(p) \\ & + \Big\{\frac{a}{1 - a^2} + z\int_0^\infty dx^2\,\rho_3(x^2)\Big\} S_F(p)(\Gamma\cdot p)\gamma_5 S_F(p) \\ & + \Big\{1 - \frac{a^2}{1 - a^2} - z\int_0^\infty dx^2\,\sigma_1(x^2)\Big\} S_F(p) \\ & - iz\int_0^\infty dx^2\, \frac{(i\Gamma\cdot p - x)\sigma_1(x^2) + \rho_2(x^2) + i(\Gamma\cdot p)\gamma_5\rho_3(x^2)}{p^2 + x^2 - i\epsilon}, \end{aligned} \tag{6}$$

where $S_F(p) = i/(i\Gamma \cdot p + m)$, spectral functions σ_1, ρ_2, and ρ_3 are all real and do not include the $\delta(x^2 - m^2)$ function, and z is a real constant. The integrals $\int_0^\infty dx^2 [(m - x)\sigma_1(x^2) + \rho_2(x^2)]$, $\int_0^\infty dx^2 \rho_3(x^2)$, and $\int_0^\infty dx^2 \sigma_1(x^2)$ diverge logarithmically in perturbation calculations. The renormalization of the modified propagator means that the unknown constants δm, z, and a are so chosen that all divergent integrals disappear from (6). This is possible because the expression (6) involves three divergent integrals and three unknown constants. When the parity-nonconserving counter term $\bar{\psi}\gamma_\mu\gamma_5(\partial/\partial x_\mu)\psi$ is not introduced in (3), the modified propagator is expressed by (6) by taking $\Gamma_\mu \to \gamma_\mu$ and a = 0. Then the expression involves three divergent integrals and two unknown constants δm and z. Thus consistent renormalization is impossible in this case. This is the reason why we introduce the parity-nonconserving counter term in (3).

We want to renormalize the wave function $\psi(p)$ for the free decorated particle with momentum p as

$$\psi(p) = z_2^{\frac{1}{2}} \psi_I(p) \tag{7a}$$

and S_F^1 as

$$S_F^1(p) = z_2 S_{F,re}^1(p) , \tag{7b}$$

where ψ_I represents the wave function in interaction representation, $S_{F,re}^1$ is the renormalized modified propagator and is divergence-free, and z_2 is a c-number constant. By the definition (7), z_2 is the bare-state probability of the decorated particle and should satisfy the relation

$$1 \geq z_2 \geq 0 . \tag{8}$$

Then the renormalization constants and the renormalized modified propagator are expressed as

$$\begin{cases} \delta m = z_2 \int_0^\infty dx^2 \left[(m - x)\, \rho_1(x^2) + \rho_2(x^2) + a m\, \rho_3(x^2) \right] , \\ z_2^{-1} = \dfrac{1 - a^2}{1 - 3a^2} \int_0^\infty dx^2 \left[\rho_1(x^2) + a\, \rho_3(x^2) \right] , \\ \dfrac{a}{1 - a^2} = - z_2 \int_0^\infty dx^2\, \rho_3(x^2), \end{cases} \tag{9}$$

and

$$S_{F,re}^{1}(p) = - i \int_0^\infty dx^2 \, \frac{(i\Gamma \cdot p - x)\, \rho_1(x^2) + \rho_2(x^2) + i\Gamma \cdot p\, \gamma_5\, \rho_3(x^2)}{p^2 + x^2 - i\epsilon} , \tag{10}$$

where $\rho_1(x^2) = \delta(x^2 - m^2) + \sigma_1(x^2)$.

When parity is conserved, as we know, after the mass and the charge renormalizations the S matrix is free from divergences. When parity is not conserved, on the other hand, the S matrix is not free from the divergent integral $\int_0^\infty dx^2\, \rho_3(x^2)$ after the usual mass and charge renormalizations. Thus the parity-nonconserving counter term is introduced. Then the divergent integral appears only in the expression for a in Γ_μ after the renormalization, as was shown by Eq. (10). Further a is a finite real constant. From Eq. (8) and the third of Eqs. (9), in fact, one obtains

$$a^2 \leqslant \frac{1}{3} . \tag{11}$$

The constant a is expressed as

$$a = \frac{1 - \left\{ 1 + \left[2 z_2 \int_0^\infty dx^2\, \rho_3(x^2) \right]^2 \right\}^{1/2}}{2 z_2 \int_0^\infty dx^2\, \rho_3(x^2)} . \tag{12}$$

One may now raise the question: Although a is a function of the divergent integral $\int_0^\infty dx^2\, \rho_3(x^2)$, can it be a finite constant? In perturbation

calculations, z_2^{-1} and $\int_0^\infty dx^2 \rho_3(x^2)$ diverge logarithmically. Thus it is possible that the combination

$$z_2 \int_0^\infty dx^2 \rho_3(x^2)$$

remains a finite constant. When the parity-nonconserving counter term is not introduced, the S matrix includes the divergent integral itself [not the combination $z_2 \int_0^\infty dx^2 \rho_3(x^2)$].

2. Pseudoscalar Charge Density

Because of the gauge invariance of our theory, even if parity is not conserved, the differential operator

$$-\frac{\partial}{\partial p_\mu}$$

is the operator to insert a photon vertex without momentum transfer into the part of the Feynman diagram which represents any charged particle with four-momentum p_μ. From the expression (10), the renormalized photon vertex without momentum transfer for any charged spin-$\frac{1}{2}$ particle on the mass shell is given by

$$\bar{\psi}(p) S_F^{-1}(p) \left[-\frac{\partial}{\partial p_\mu} S_{F,re}^{1}(p) \right] S_F^{-1} \psi(p) = \bar{\psi}(p) \Gamma_\mu \psi(p) . \qquad (13)$$

This equation means that the charge density of a free charged particle is given by

$$\rho(x) = e \bar{\psi}(x) \beta (1 + a\gamma_5)(1 - a^2)^{-1/2} \psi(x) = e \psi^*(1 + a\gamma_5)(1 - a^2)^{-1/2} \psi, \qquad (14)$$

which consists of scalar part $\psi^*(1 - a^2)^{-1/2} \psi$ and pseudoscalar part $\psi^* a\gamma_5 (1 - a^2)^{-1/2} \psi$. We shall call the latter the pseudoscalar charge

density. It should be stressed that the pseudoscalar charge density is not arbitrarily introduced in our theory but is induced by parity-nonconserving interactions.

Under the requirements of covariance, gauge invariance, and CP invariance, the most general form of the photon vertex on the mass shell for a spin-$\frac{1}{2}$ particle is expressed as

$$\bar{\psi}(p_1)\, j_\mu(p_1, p_2)\, \psi(p_2)$$

$$= i e \bar{\psi}(p_1) \{ \Gamma_\mu F_1(q^2) + \frac{1}{2m} \sigma_{\mu\nu} q_\nu F_2(q^2)$$

$$+ \frac{1}{m^2} [q^2 \Gamma_\mu - (\Gamma \cdot q) q_\mu] \gamma_5 F_3(q^2) \} \psi(p_2), \qquad (15)$$

where $q_\mu = (p_1 - p_2)_\mu$ and all form factors F_i ($i = 1, 2,$ or 3) are all real when q_μ is a space-like vector. The renormalization condition (13) leads to the conditions

$$F_1(0) = \begin{cases} 1 & \text{for charged particles ,} \\ 0 & \text{for neutral particles .} \end{cases} \qquad (16)$$

The first term in Eq. (15) consists of a vector part $(1 - a^2)^{-1/2} F_1$ and an induced pseudovector part $a\gamma_5 (1 - a^2)^{-1/2} F_1$ which relates to the pseudoscalar charge density. Thus both charged and neutral spin-$\frac{1}{2}$ particles have induced pseudoscalar charge densities. The second and third terms in (15) are moment and anapole-moment terms, respectively. In the next section we shall show that the induced pseudoscalar charge density has physical meaning when the particle is not free.

3. Observability of the Pseudoscalar Charge Density

Since two gamma matrices satisfy the same commutation relations (2) and (4), we have the theorem: There exists a nonsingular matrix S such that

$$\Gamma_\mu = S \gamma_\mu S^{-1} \tag{17}$$

and S is unique except for an arbitrary multiplicative factor. In fact the matrix S is given by

$$S = \left[\frac{(1 - a^2)^{1/2}}{2[1 + (1 - a^2)^{1/2}]} \right]^{\frac{1}{2}} \left[1 + \frac{(1 - a\gamma_5)}{(1 - a^2)^{1/2}} \right] . \tag{18}$$

Starting from the Lagrangian density

$$L = - \bar{\psi}'(x) \left[\gamma_\mu \frac{\partial}{\partial x_\mu} + m \right] \psi(x) , \tag{19}$$

the transformation S leads to

$$L = - \bar{\psi}(x) \left[\Gamma_\mu \frac{\partial}{\partial x_\mu} + m \right] \psi(x), \tag{20}$$

where $\psi = S\psi'$ and $\bar{\psi} = \psi^* \beta$. Since two gamma matrices γ_μ and Γ_μ satisfy the same commutation relation, as we know, two Lagrangian densities (19) and (20) describe the same free spin-$\frac{1}{2}$ particle. Thus the constant a has no physical meaning when the particle is free.

We shall show in the following that, when the spin-$\frac{1}{2}$ particle interacts with other fields, a can have a physical meaning. Even if a has no physical meaning for free particles, this is possible because the transformation matrix S is not unitary, i.e., because

$$S^{*}S = (1 - a\gamma_5)(1 - a^2)^{-1/2} \neq 1 \qquad \text{for } a \neq 0 . \tag{21}$$

We want to calculate the energy of a charged spin-$\frac{1}{2}$ particle in a weak external electromagnetic field. For simplicity, the vertex operator (15) is approximated by taking $F_1 \equiv 1$, $F_2 = F_3 \equiv 0$. Then the Lagrangian density of the system is given by

$$L' = -\bar{\psi}' \left[\gamma_\mu \left(\frac{\partial}{\partial x_\mu} - ieA_\mu \right) + m \right] \psi' \tag{22a}$$

and

$$L = -\bar{\psi} \left[\Gamma_\mu \left(\frac{\partial}{\partial x_\mu} - ieA_\mu \right) + m \right] \psi , \tag{22b}$$

which are obtained from (19) and (20), respectively. From these Lagrangian densities, one obtains the equations of motion for two component wave functions ψ_1' and ψ_1, namely

$$E'\psi_1' = \{ m + e\phi + \frac{1}{2m} \vec{\sigma}(\vec{p} - e\vec{A}) \cdot \vec{\sigma}(\vec{p} - e\vec{A}) \} \psi_1' + 0\left(\frac{1}{m^2}\right) \tag{23}$$

and

$$\begin{aligned} E\psi_1 = \Bigg\{ & m + e\phi + \frac{1}{2m} \vec{\sigma}(\vec{p} - e\vec{A}) \cdot \vec{\sigma}(\vec{p} - e\vec{A}) \\ & - \frac{1}{2m[\,1 + (1 - a^2)^{1/2}\,]} [\,(a\, e\phi)\, \vec{\sigma}(\vec{p} - e\vec{A}) + \vec{\sigma}(\vec{p} - e\vec{A})(a\, e\phi)\,] \\ & + \frac{1}{2m[\,1 + (1 - a^2)^{1/2}\,]^2} (a\, e\phi)^2 + 0\left(\frac{1}{m^2}\right) \Bigg\} \psi_1 , \end{aligned} \tag{24}$$

where

$$\psi' = \begin{pmatrix} \psi_1' \\ \psi_2' \end{pmatrix} , \qquad \psi = \begin{pmatrix} \psi_1 \\ \psi_2 \end{pmatrix} ,$$

and $A_\mu = (\vec{A}, \phi)$. The extra terms that appear in Eq. (24) but not in (23) are proportional to $(a\, e\phi)$ which comes from the interaction of the pseudoscalar charge density $a\, e\gamma_5 (1 - a^2)^{-1/2}$ with the external field ϕ. Thus it has been shown that the pseudoscalar charge density is an observable. It should be noted that our pseudoscalar charge density is not obtained by the nonunitary transformation (17) but it is induced by parity-nonconserving interactions as was discussed in Secs. 2 and 3.

So far we have considered only renormalizable interactions and have proved the existence of the observable pseudoscalar charge density. When unrenormalizable interactions are taken into account, we cannot prove the existence because unrenormalizable theory is a phenomenological theory and it is not a consistent theory. However, we may argue that unrenormalizable and parity-nonconserving interactions would be one of the sources of the pseudoscalar charge density.

V-53-1. SYMMETRIES OF ELEMENTARY PARTICLES

Harry J. Lipkin and Amnon Katz (51151-01)

AN INVESTIGATION OF ω-ϕ MIXING

So far only one pair of elementary-particle resonances has been found to have the same quantum numbers (spin, parity, and isotopic spin), namely the ω and ϕ vector-meson resonances. Since all their quantum numbers are the same, one might expect them to have a similar behavior since any mechanism allowed for production and decay of one is also allowed for the other. Experimentally, this turns out not to be the case. The ω is produced copiously in pion-nucleon reactions whereas the ϕ is not. The ω decays into three pions whereas the corresponding decay mode of the ϕ is either absent or very strongly suppressed.

The unitary symmetry theory classifies particles into multiplets and requires one particle having the quantum numbers of the ω and the ϕ to be in an octet together with the K* and ρ vector mesons while the other must be by itself in a singlet. Sakurai has pointed out that a consistent description in agreement with the unitary-symmetry mass formula is possible if one assumes that the singlet and octet states are mixed and that the physical ω and physical ϕ mesons are both linear combinations of the unitary singlet and the unitary octet with roughly equal amplitudes. This makes the difference in the behavior of the ϕ and the ω even more mysterious. In order to explain the lack of coupling between the physical ϕ meson and pions, it is necessary to assume that both the unitary singlet and the unitary octet components are coupled approximately equally in these pion reactions and that the singlet and octet contributions cancel one another out for reactions involving the ϕ and add together constructively in reactions involving the ω.

The purpose of this work was to investigate whether a single dynamical model could explain both the $\omega\phi$ mixing and the weak coupling of the ϕ in pion reactions. The basic idea is that the unitary singlet and the unitary octet are initially approximately degenerate and that the mixing is predominantly produced by a virtual transition of the vector meson into a ρ and a π and back again. This mechanism produces a mixing of the unitary singlet and the unitary octet whose eigenstates are obtained by diagonalizing a two-by-two matrix. One of the eigenstates turns out to be automatically decoupled from the $\rho\pi$ channel.

A more complete calculation is presented in which other two-meson and baryon-antibaryon virtual states are considered. The couplings are assumed to be consistent with unitary symmetry with the symmetry broken by the use of experimental values for the masses of the particles occurring in the intermediate states. Although a straightforward calculation leads to divergent integrals which cannot be evaluated, the values of certain functions of these integrals can be determined by fitting experimental mass splittings of the vector mesons. The result of the calculation is a very small coupling of the ϕ to the $\rho\pi$ channel, a reasonably good fit to the vector-meson masses (there are four mass differences and three free parameters so that the fit is significant), and a value (in reasonable agreement with experiment) for the ratio of the width of the ω to the width of the ρ.

PUBLICATIONS SINCE THE LAST REPORT

PAPERS

DETERMINATION OF ISOTOPIC DIFFERENCES IN MESIC X-RAYS IN CALCIUM
H. L. Anderson,* C. S. Johnson,* E. P. Hincks,† S. Raboy, and C. C. Trail . . . . . (Project I-57)
Phys. Letters 6, 261-264 (15 September 1963)

EQUILIBRIUM COMPOSITION OF SULFUR VAPOR
J. Berkowitz and J. R. Marquart . . . . . (Project II-29)
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L. M. Bollinger, R. E. Coté, R. T. Carpenter, and J. P. Marion . . . . . (Project I-7)
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W. J. Childs, L. S. Goodman, and D. von Ehrenstein . . . . . (Project I-80)
Phys. Rev. 132, 2128-2135 (1 December 1963)

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M. N. Hack . . . . . (Project V-26)
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THEORY OF SPIN-$\frac{1}{2}$ PARTICLES WITH PARITY-NONCONSERVING INTERACTIONS
K. Hiida . . . . . (Project V-47)
Phys. Rev. 132, 1239-1248 (1 November 1963)

PROPERTIES OF CAPTURE GAMMA-RAY SPECTRA OF s-WAVE AND p-WAVE RESONANCES IN Zr, Nb, AND Mo
H. E. Jackson . . . . . (Project I-7)
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H. E. Jackson . . . . . (Project I-7)
Phys. Rev. Letters 11, 378-380 (15 October 1963).
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A DYNAMICAL MODEL FOR ω-φ MIXING
A. Katz and H. J. Lipkin . . . . . (Project V-53)
Phys. Letters 7, 44-47 (15 October 1963)

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†National Research Council of Canada, Ottawa, Ontario, Canada.

UNITARY SYMMETRY IN PHOTOPRODUCTION AND OTHER ELECTROMAGNETIC INTERACTIONS
C. A. Levinson,* H. J. Lipkin, and S. Meshkov† . . .(Project V-53)
Phys. Letters 7, 81-84 (15 October 1963)

ON THE PRODUCTION OF THE Z^- (Ω^-) BARYON
H. J. Lipkin, C. A. Levinson,* and S. Meshkov† . . .(Project V-53)
Phys. Letters 7, 159-161 (1 November 1963)

BOSON COUPLINGS AND HIGHER SYMMETRIES
H. J. Lipkin . . . . . .(Project V-53)
Phys. Letters 7, 221-222 (15 November 1963)

DISSOCIATION ENERGIES OF SOME METAL SULFIDES
J. R. Marquart and J. Berkowitz . . . . . (Project II-28)
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STEREOSCOPIC TELEVISION SYSTEM
W. J. O'Neill . . . . . (Project I-22)
Nucl. Instr. and Methods 24, 45-48 (July 1963)

FLUXOID QUANTIZATION, PAIR SYMMETRY, AND THE GAP ENERGY IN THE CURRENT-CARRYING BARDEEN-COOPER-SCHRIEFFER STATE
M. Peshkin . . . . . (Project V-33)
Phys. Rev. 132, 14-20 (1 October 1963)

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N. Rosenzweig . . . . . .(Project I-15)
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J. P. Schiffer . . . . . (Project I-30)
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THE PARITY OF THE NEUTRAL K MESON
R. Spitzer . . . . . .(Project V-39)
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*Weizmann Institute of Science, Rehovoth, Israel.

†National Bureau of Standards, Washington, D.C.

REFLECTION OF A PLANE SOUND WAVE FROM A SINUSOIDAL SURFACE

J. L. Uretsky . . . . . (Project V-50)
J. Acoust. Soc. Am. 35, 1293-1294 (1963)

THE OBSERVATION OF Be^8 IN THE GROUND STATE

J. A. Weinman and R. K. Smither . . . . . (Project I-21)
Nucl. Phys. 45, 260-272 (1963)

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SOME SIMPLE FEATURES OF THE MÖSSBAUER EFFECT

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POLARIZED SPECTRA AND HYPERFIND STRUCTURE IN Fe^{57}

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L. Meyer-Schützmeister, R. S. Preston, and S. S. Hanna . . . . . (Project I-19)
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RECOILLESS RESONANCE ABSORPTION OF GAMMA RAYS IN Fe^{57}
J. P. Schiffer and W. Marshall . . . . . .(Project I-19)
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THE ARGONNE 3-PARAMETER ANALYZER

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Proceedings of the Conference on Utilization of Multi-parameter Analyzers in Nuclear Physics, Grossinger's, New York, 1962 (U. S. Atomic Energy Commission Report NYO 10595), p. 59

THERMODYNAMICS AND KINETICS OF VAPORIZATION PROCESSES STUDIED BY MASS SPECTROMETRY

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J. H. Roberts, P. C. Chow, and G. R. Ringo. . .(Project I-132)
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A. M. Saperstein and J. L. Uretsky . . . . . .(Project V-49)
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(d,t) REACTIONS ON NUCLEI NEAR N = 28

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SENIORITY MIXING IN $1f_{7/2}$ NUCLEI
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THE BOOTSTRAP METHOD IN THE PHYSICS OF STRONG INTERACTIONS
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POPULATION OF THE FIRST T = 1 STATE IN B^{10} BY THE REACTION $C^{12}(d,\alpha)B^{10}$
R. G. Allas, J. R. Erskine, L. Meyer-Schützmeister, and D. von Ehrenstein . . . . . .(Project I-21)
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CAPTURE REACTIONS $Al^{27}(p,\gamma)Si^{28}, ^{28*}$ IN THE GIANT-DIPOLE REGION OF Si^{28}
R. G. Allas, S. S. Hanna, L. Meyer-Schützmeister, R. E. Segel, and P. P. Singh . . . . . .(Project I-21)
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ISOTOPE SHIFT IN THE K_α TRANSITION OF μ-MESIC ATOMS OF Ca^{40} AND Ca^{44}
H. L. Anderson, C. S. Johnson, E. P. Hincks, S. Raboy, and C. C. Trail . . . . . .(Project I-57)
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LEVELS IN Ga^{66} FROM THE DECAY OF Ge^{66}
H. H. Bolotin . . . . . (Project I-6)
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ACCURACY OF THE QUASIPARTICLE RANDOM-PHASE APPROXIMATION (QPRPA) IN NUCLEAR-STRUCTURE CALCULATIONS
S. Cohen, R. D. Lawson, M. H. Macfarlane, and M. Soga . . . . . .(Project V-9)
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POLARIZATION OF NEUTRONS SCATTERED FROM F, Na, Al, AND P
A. J. Elwyn, R. O. Lane, and A. Langsdorf, Jr. . . . .(Project I-18)
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D. S. Gemmell, L. L. Lee, Jr., A. Marinov, and J. P. Schiffer . . . . . .(Project I-27)
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FLUCTUATIONS IN PARTIAL RADIATION WIDTHS OF U^{239}
H. E. Jackson . . . . . .(Project I-7)
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Chicago meeting of the American Physical Society (cont'd.)

DYNAMICAL ASPECTS OF MESON-BARYON SCATTERING IN THE OCTET MODEL
A. W. Martin and K. C. Wali . . . . . (Project V-52)
Bull. Am. Phys. Soc. 8, 515 (18-19 October 1963)

OPTICAL-MODEL ANALYSIS OF NEUTRON SCATTERING FROM NUCLEI NEAR A = 100 AT LOW ENERGIES
J. E. Monahan, A. J. Elwyn, R. O. Lane, and A. Langsdorf, Jr. . . . . (Project I-18)
Bull. Am. Phys. Soc. 8, 513 (18-19 October 1963)

SPIN-PARITY ANALYSIS AT ARBITRARY PRODUCTION ANGLE
M. Peshkin . . . . . (Project V-5)
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SEARCH FOR BOUND STATES IN He^4
J. L. Yntema, S. S. Hanna, and R. E. Segel. . . . (Project I-22)
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ANL TOPICAL REPORT

Proceedings of the International Conference on Nuclear Physics with Reactor Neutrons, 15-17 October 1963, edited by F. E. Throw, Argonne National Laboratory Topical Report ANL-6797 (490 pages)

STUDENT REPORTS

ISOMER SHIFT IN VANADIUM-ALUMINUM ALLOYS AND APPLICATION OF THE MÖSSBAUER EFFECT TO METEORITE ANALYSIS
D. J. Bailey . . . . . (Project I-19)
ACM student report to Beloit College (August 1963)

GRAPHING OF THEORETICAL VALUES OF THE COEFFICIENTS OF LEGENDRE POLYNOMIALS
B. Blumenstein . . . . . (Project I-18)
Co-op student report to the University of Florida (August 1963)

CO-OPERATIVE EDUCATION WORK REPORT
G. Hoffman . . . . . (Project I-80)
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CALCULATION AND ANALYSIS OF ELASTIC SCATTERING CROSS SECTIONS BY MEANS OF COMPUTER PROGRAMS

D. Howe . . . . . (Project I-18)
Co-op student report to the University of Detroit (November 11, 1963)

IONIC IMPACT PHENOMENA ON METAL SURFACES IN THE RUTHERFORD COLLISION REGION

T. Klopcic . . . . . (Project II-23)
ACM student report to Knox College (August 1963)

THE "SUPER" PROGRAM

T. Lawler . . . . . (Project I-31)
Co-op student report to Marquette University (September 1963)

METHODS OF NUCLEAR RESEARCH

R. E. Malmin . . . . . (Project I-21)
Co-op student report to the University of Detroit (November 14, 1963)

A STUDENT'S DISCUSSION OF THE NUCLEAR OPTICAL MODEL

D. J. Mueller
Co-op student report to Northwestern University (September 9, 1963)

THE $C^{12}(d,\alpha)B^{10}$ AND $K^{39}(p,\alpha)Ar^{36}$ REACTIONS

C. D. Pruett . . . . . (Project I-21)
Co-op student report to Purdue University (September 1963)

THE MÖSSBAUER EFFECT IN Sn^{119} AND Mn_3Sn

M. P. Schulhof . . . . . (Project I-19)
ACM student report to Grinnell College (August 1963)

EFFICIENCY OF SCINTILLATION DETECTORS FOR POLARIZATION EXPERIMENTS

C. A. Walker . . . . . (Project I-18)
Co-op student report to the University of Florida (December 20, 1963)

INFORMAL REPORT

UNITARY SYMMETRY FOR PEDESTRIANS (U-SPIN, I-SPIN, V-ALL SPIN FOR I-SPIN)

H. J. Lipkin . . . . . (Project V-53)
Informal report

ADDITIONAL PAPERS ACCEPTED FOR PUBLICATION

VAPORIZATION PROCESSES INVOLVING SULFUR
J. Berkowitz and W. A. Chupka . . . . . .(Project II-29)
J. Chem. Phys. (15 January 1964)

p-WAVE RESONANCES OF Th^{232}
L. M. Bollinger and G. E. Thomas . . . . . .(Project I-3)
Phys. Letters (15 December 1963)

NOTE ON A CRUDE NUCLEON-NUCLEON POTENTIAL
F. Coester and E. Yen (University of Iowa). . . .(Project V-19)
Nuovo Cimento (December 1963)

NONLOCAL OPTICAL-MODEL ANALYSIS OF NEUTRON SCATTERING FROM NUCLEI NEAR A = 100 AT ENERGIES BELOW 1 MEV
A. J. Elwyn, R. O. Lane, A. Langsdorf, Jr., and
J. E. Monahan . . . . . (Project I-18)
Phys. Rev. (13 January 1964)

THE $Ta^{181}(d,p)Ta^{182}$ REACTION AND A BAND-MIXING ANALYSIS OF THE OBSERVED ENERGY LEVELS
J. R. Erskine and W. W. Buechner . . . . . .(Project I-23)
Phys. Rev. (13 January 1964)

NUCLEAR LEVELS OF F^{20}
C. T. Hibdon . . . . . .(Project I-98)
Phys. Rev. (27 January 1964)

POLARIZATION AND DIFFERENTIAL CROSS SECTION OF NEUTRONS SCATTERED FROM Be^{9} : PARITIES OF THE 7.37- AND 7.54-MEV STATES IN Be^{10}
R. O. Lane, A. J. Elwyn, and A. Langsdorf, Jr. . . .(Project I-18)
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MESON-BARYON RESONANCES IN THE OCTET MODEL
A. W. Martin and K. C. Wali . . . . . .(Project V-52)
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NEUTRON PRODUCTION BY 450-MEV PROTONS
J. W. Meadows, G. R. Ringo, and A. B. Smith. . .(Unattached)
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AN S-STATE INTERACTION WITH PAIRING-TYPE PROPERTIES IN THE NUCLEAR SHELL MODEL
S. P. Pandya . . . . . (Project V-6)
Phys. Letters (15 December 1963)

SPIN AND PARITY ANALYSIS AT ALL PRODUCTION ANGLES
M. Peshkin . . . . . (Project V-5)
Phys. Rev. (27 January 1964)

FAIRLY LOW-ENERGY PION-PION SCATTERING II
A. M. Saperstein and J. L. Uretsky . . . . . .(Project V-49)
Phys. Rev.

SLOW-NEUTRON-CAPTURE GAMMA RAYS IN Hg^{200}
R. E. Segel, R. K. Smither, and R. T. Carpenter. . .(Project I-7)
Phys. Rev. (10 February 1964)

VECTOR MESON DECAYS IN UNITARY SYMMETRY
K. Tanaka . . . . . .(Project I-45)
Phys. Rev.

STATES IN Si^{28} WITH $12.7 < E_x < 13.7$ MEV BY (α,γ) AND (α,α) REACTIONS ON Mg^{24}
J. A. Weinman, L. Meyer-Schützmeister, and
L. L. Lee, Jr. . . . . .(Project I-21)
Phys. Rev. (10 February 1964)

PERSONNEL CHANGES IN THE ANL PHYSICS DIVISION

NEW MEMBERS OF THE DIVISION

Resident Research Associates

Dr. Arnold Bodmer, Department of Theoretical Physics, Manchester University, England. Hypernuclei. Came to Argonne on 18 October 1963. (Host: Murray Peshkin.)

Dr. Harnam S. Hans, Department of Physics, Texas A and M University, College Station, Texas. Resonance scattering of neutron-capture gamma rays. Came to Argonne on 16 September 1963. (Host: L. M. Bollinger.)

Dr. John K. Perring, Principal Scientific Officer, A. E. R. E., Harwell, England. Nucleon-nucleon interaction. Came to Argonne on 9 September 1963. (Host: M. Peshkin.)

Resident Research Associates (Post-Doctoral)

Dr. Robert W. Childers, Department of Physics, Vanderbilt University. Problems in particle physics. Came to Argonne on 5 September 1963. (Host: M. Peshkin.)

Dr. Arthur I. Namenson, Department of Physics, Columbia University. (n, γ) reactions and level schemes. Came to Argonne on 14 October 1963. (Host: R. K. Smither.)

Dr. Rolf H. Siemssen, II. Institut für Experimentalphysik, University of Hamburg, Germany. Nuclear spectroscopic studies of light and medium-weight nuclei with stripping and pickup reactions. Came to Argonne on 2 October 1963. (Host: L. L. Lee, Jr.)

Dr. Zeev Vager, Physics Division, Weizmann Institute of Science, Israel. Giant resonance of Ca^{40} at the tandem Van de Graaff. Came to Argonne on 18 November 1963. (Hosts: L. Meyer-Schützmeister and R. E. Segel.)

Student Aides

Mr. Edward Gill, Kalamazoo College. Working with W. A. Chupka and J. Berkowitz on electronic, mechanical, optical, and photographic studies of the phenomena involved in the jet produced when light of very high power density (from a ruby laser) is focused on a graphite surface. Came to ANL on 1 October 1963.

Student Aides (ACM)

Mr. Frederick R. Berg, Grinnell College, Iowa. Working with J. Heberle on the nuclear Zeeman effect in Sn^{119} in Nb_3Sn. Came to ANL on 3 September 1963.

Mr. George Dinolt, Lawrence College, Wisconsin. Working with S. B. Burson on the decay schemes of short-lived radionuclides. Came to ANL on 3 September 1963.

Mr. William Elliott, Ripon College, Wisconsin. Working with J. P. Schiffer on charge distributions of nuclei. Came to ANL on 3 September 1963.

Co-op Technicians

Mr. Rudolph Henninger, Northwestern University. Working with J. P. Schiffer on analysis of data on charged-particle reactions. Came to ANL on 23 December 1963.

Mr. Denis Howe, University of Detroit. Working with L. L. Lee on analysis of data on charged-particle reactions. Came to ANL on 11 September 1963.

Mr. Ronald Malmin, University of Detroit. Working with Luise Meyer-Schützmeister on the population of the first T = 1 level in B^{10} by the $C^{12}(d, \alpha) B^{10}$ reaction at the tandem Van de Graaff. Came to ANL on 3 September 1963.

Mr. Kenneth Swenson, University of Michigan. Working with M. Kaminsky on pulsed-molecular-beam mass spectrometry. Came to ANL on 16 October 1963.

Miss Claudia Walker, University of Florida. Working with R. O. Lane and A. J. Elwyn on an experiment employing scintillation detectors for neutron polarization experiments. Came to ANL on 3 September 1963.

Technicians

Mr. Milton Feryance joined the Physics Division on 18 November 1963 as a Research Technician (Junior) to work with R. Amrein.

Mr. Vladimir W. Nemec joined the Physics Division on 14 October 1963 as a Research Technician (Junior) to work with William Evans.

Mr. James Smith joined the Physics Division on 26 September 1963 as a Research Technician to work with D. S. Gemmell.

Secretaries

Miss Barbara Baldwin joined the Physics Division on 11 November 1963 as secretary in F wing.

Mrs. Jeanette Vendel joined the Physics Division on 11 November 1963 as secretary in B wing.

Data Analyzer

Miss Paulette Radd joined the Physics Division on 16 September 1963 to assist the users of the magnetic spectrograph in extracting data from nuclear track plates.

LEAVES OF ABSENCE

Dr. Benjamin Zeidman left ANL on 1 October 1963 on a fellowship for a year of experimental and theoretical investigations in nuclear structure at the Institute for Theoretical Physics, Blegdamsvej 17, Copenhagen, Denmark. He expects to return to Argonne on 20 September 1964.

DEPARTURES

Mrs. Carol Lindgren, who has been the receptionist of the Physics Division since 10 July 1961, terminated on 13 November 1963.

Dr. Harry J. Lipkin, who has been a Resident Research Associate in the Physics Division since 3 June 1963, terminated on 1 November 1963 to return to the Weizmann Institute of Science, Rehovoth, Israel.

Mr. Marvin C. Mertz, who has been a Resident Student Associate (Thesis) in the Physics Division since 28 June 1961, terminated on 30 September 1963 to go to Northern Illinois University, DeKalb, Illinois.

Dr. Alvin M. Saperstein, who has been a Resident Research Associate in the Physics Division since 18 September 1962, terminated on 12 September 1963 to go to Wayne State University, Detroit, Michigan.

Mrs. JoAnn Sullivan, who has been the secretary of the theoretical wing in the Physics Division since 31 October 1960, terminated on 15 October 1963.

Dr. Bhalchandra M. Udgaonkar, who has been a Resident Research Associate in the Physics Division since 22 April 1963, terminated on 2 August 1963 to return to the Tata Institute of Fundamental Research, Bombay, India.